

Physics of Self-Interacting Electroweak Bosons *

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Abstract

I will argue why and how it is that precise measurements of the self-couplings of the weak vector bosons are a vista on the mechanism of symmetry breaking. Guided by what we have learnt from the present precision data, it is suggested which of the many so-called anomalous self-couplings should be given priority in future searches. Expected limits from the upcoming colliders on the parameters describing non minimal couplings are updated. I will also point at the complementarity between the LHC and the Next Linear Collider as concerns W physics and discuss some of the important issues about radiative corrections and backgrounds that need further studies in order that one conducts high precision analysis at high energies.

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1 Symmetry breaking and anomalous gauge bosons couplings

1.1 The mass Connection

All data to date, crowned by the results of LEP 1 (dedicated physics with 1 weak boson) have left no doubt that the Standard Model, \mathcal{SM} , has passed with flying colours all the low energy tests, even and especially at the quantum level. Yet, despite the absence of the slightest hint of any anomaly, the model has still not been elevated to the status of a fully-fledged theory. The reason for this status is essentially due to the sector in the model that implements the mass generation and the mechanism of symmetry breaking, \mathcal{SB} . It is in this sector that originates the remaining missing particle of the model, the Higgs, about which even the very precise data give no direct unambiguous clue. Add to this that an elementary scalar is unnatural, it is no wonder that almost all the *beyond the \mathcal{SM}* activity covered by the various talks at this conference is an investigation or a modification of this sector.

Whatever the structure and the particle content of this sector, we know, at least, that it contains

$$\mathcal{L}_M = M_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2} M_Z^2 Z_\mu Z^\mu \quad (1)$$

For the fermionic mass terms our knowledge is even more limited as we do not have access to all the elements of the mass matrices of the ups and downs.

The mass term (1) is the most trivial term that may be regarded as describing a self-coupling between the W 's. These couplings tell another story than the self-couplings that are present in any unbroken gauge theory like QCD, say, which originate from the kinetic part of the spin-1 boson and which describe the propagation and interaction of transverse states. In the electroweak model[‡]

$$\mathcal{L}_G = -\frac{1}{2} [Tr(\mathbf{W}_{\mu\nu} \mathbf{W}^{\mu\nu}) + Tr(\mathbf{B}_{\mu\nu} \mathbf{B}^{\mu\nu})] \quad (2)$$

These interactions only involve the field strength and are thus explicitly gauge invariant. The mass terms, that introduce the longitudinal degrees of freedom, would seem to

[‡]The conventions and definitions of the fields and matrices that I am using here are the same as those in [1].

break this crucial local gauge symmetry. The important point, as you know, is that the symmetry is not broken but rather hidden. Upon introducing auxiliary fields with the appropriate gauge transformations, we can rewrite the mass term in a manifestly local gauge invariant way (through the use of covariant derivatives). In the minimal standard model this is done through a doublet of scalars, Φ , of which one is the *physical* Higgs. The simplest choice of the doublet implements an extra global *custodial* $SU(2)$ symmetry that gives the well established $\rho = \frac{M_W^2}{M_Z^2 c_W^2} \simeq 1$:

$$\mathcal{L}_{H,M} = (\mathcal{D}_\mu \Phi)^\dagger (\mathcal{D}^\mu \Phi) - \lambda \left[\Phi^\dagger \Phi - \frac{\mu^2}{2\lambda} \right]^2 \quad (3)$$

In the case where the Higgs does not exist or is too heavy, one can modify this prescription such that only the Goldstone Bosons $\omega_{1,2,3}$, grouped in the matrix Σ , are eaten (see for instance[2]):

$$\mathcal{L}_M = \frac{v^2}{4} \text{Tr}(\mathcal{D}^\mu \Sigma^\dagger \mathcal{D}_\mu \Sigma) ; \quad \Sigma = \exp\left(\frac{i\omega_\alpha \tau^\alpha}{v}\right) \quad (v = 246 GeV) \quad (4)$$

In this so-called non-linear realisation of \mathcal{SB} the mass term (1) is formally recovered by going to the physical “frame” (gauge) where all Goldstones disappear, *i.e.*, $\Sigma \rightarrow \mathbf{1}$.

The above operators that describe the self-interaction of the vector bosons constitute the *minimal* set of operators that can be written given the well-confirmed symmetries of the weak interaction and the known content of the \mathcal{SM} spectrum. In this sense, the non-linear realisation is even more economical since it does not appeal to the still missing Higgs. These operators are minimal not only in the sense of their fields content but also in the sense that these are the lowest dimension operators that we can write. In the case of the non-linear realisation it is more appropriate to talk about operators with the least number of derivatives. One expects that, in the absence of a direct observation of new particles especially those that emerge from the mass sector, phenomena related to \mathcal{SB} can be described in terms of higher order terms constructed in the mould of (3,4). These induce new weak bosons self-couplings. To investigate their presence we would then study interactions involving *longitudinal* vector bosons.

Of course, one can construct other operators describing vector bosons self-couplings on the mould of the *universal* kinetic term (2) which is explicitly gauge invariant. In this case it is worth keeping in mind that these types of anomalies will not be telling us much about symmetry breaking, but only that there may be some weakly heavy interacting particles. The oldest example of such operators for the transverse modes is the celebrated

Euler-Heisenberg Lagrangian that describes (in the first order) an *anomalous* 4-photon coupling. Given my bias about the importance of effects intimately related to \mathcal{SB} I will not be concentrating much on this type of anomalies, this is the first level where I would like to discriminate between origins of anomalies.

Another example of an effective Lagrangian that has proved more revealing and rich in physics is the effective chiral Lagrangian that describes the interaction of pions out of which one has learnt so much about the interaction of hadrons. Likewise, one hopes that the electroweak equivalent (generalisation of (4) where the pions are to be identified with the pseudo-Goldstone bosons) will teach us something about symmetry breaking. I will also take the biased point of view (level 2 of discrimination) that if one still pursues the description of anomalous couplings within the light Higgs linear approach, then it may be more educating to probe the characteristics and the couplings of the Higgs. But this is not the subject of my talk.

To summarise at this point, the type of self-couplings that, in my view, deserve the highest priority are those that one has to probe in the eventuality that there is no Higgs. This is because I consider that if the Higgs is light one has already learnt a great deal about \mathcal{SB} , that the weak interaction will remain weak at TeV energies and that one should probably concentrate on studying the spectra of the New Physics that is associated with the symmetry that naturally accommodates a light Higgs, SUSY.

1.2 Phenomenological Parameterisation

The purpose of my rather long introduction was to stress the connection between the investigations of the \mathcal{SB} sector through the study of anomalous gauge bosons couplings. These would be parameterised by operators of higher dimensions or of higher order in the energy expansion than those in (2 - 4). This, of course, is suggestive of an ordering of operators with respect to the scale of the new physics (the \mathcal{SB} scale in the point of view I am taking). Of course, if one is doing experiments at an energy near this scale this ordering makes no sense and one should consider all the tower of operators. In this case, especially for the longitudinal gauge bosons, our underlying gauge symmetry principle that is instrumental for the ranking, such as to filter only a small subset of operators, would not be of much help. Still, we can use the *exact non-broken* symmetries like the $U(1)_{\text{QED}}$ and Lorentz invariance to write all the possible operators that can give an effect to a particular situation. The *phenomenological* parameterisation of the $WW\gamma$ and WWZ vertex of HPZH [3] has been written for the purpose of studying $e^+e^- \rightarrow W^+W^-$, the

bread-and-butter of LEP2. The same parameterisation, although as general as it can be for $e^+e^- \rightarrow W^+W^-$, may not be necessarily correct nor general when applied to other situations. In principle, if one is guiding by this general principle of keeping only the NON-BROKEN symmetries, one should write a new set of operators for every new situation. This does not necessarily contain all the operators of HPZH. This is one of the shortcomings. Nonetheless, the HPZH parameterisation has become popular enough in discussing anomalies that I will refer to it quite often as a common ground when comparing various approaches and “data”. To keep the discussion tractable (lack of time) I will only pick out the \mathcal{C} and \mathcal{P} conserving parts of this parameterisation otherwise one has to consider in all generality 13 couplings. Indeed it has been shown[4] that a particle of spin-J which is not its own anti-particle can have, at most, $(6J + 1)$ electromagnetic form-factors including \mathcal{C} , \mathcal{P} and \mathcal{CP} violating terms. The same argument tells us [4] that if the “scalar”-part of a massive spin-1 particle does not contribute, as is the case for the Z in $e^+e^- \rightarrow W^+W^-$, then there is also the same number of invariant form-factors for the spin-1 coupling to a charged spin-J particle. The \mathcal{C} and \mathcal{P} conserving part of the HPZH[3] parameterisation is

$$\begin{aligned} \mathcal{L}_{WWV} = & -ie \left\{ \left[A_\mu \left(W^{-\mu\nu} W_\nu^+ - W^{+\mu\nu} W_\nu^- \right) + \overbrace{(1 + \Delta\kappa_\gamma)}^{\kappa_\gamma} F_{\mu\nu} W^{+\mu} W^{-\nu} \right] \right. \\ & + \cotg\theta_w \left[\overbrace{(1 + \Delta\mathbf{g}_1^Z)}^{g_1^Z} Z_\mu \left(W^{-\mu\nu} W_\nu^+ - W^{+\mu\nu} W_\nu^- \right) + \overbrace{(1 + \Delta\kappa_Z)}^{\kappa_Z} Z_{\mu\nu} W^{+\mu} W^{-\nu} \right] \\ & \left. + \frac{1}{M_W^2} (\boldsymbol{\lambda}_\gamma F^{\nu\lambda} + \boldsymbol{\lambda}_Z \cotg\theta_w Z^{\nu\lambda}) W_{\lambda\mu}^+ W_{\nu}^{-\mu} \right\} \end{aligned} \quad (5)$$

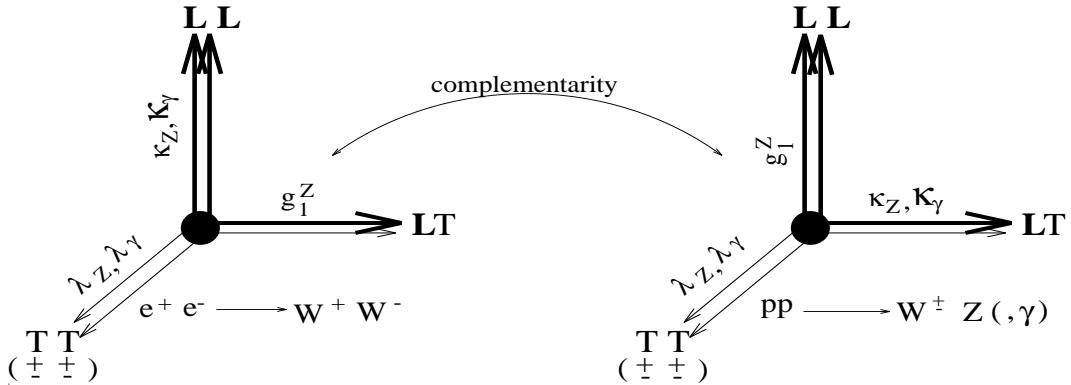
For those not working in the field and who want to get a feeling for what these form factors mean, suffice it to say that the combination $\mu_W = e(2 + \Delta\kappa_\gamma + \lambda_\gamma)/2M_W$ describes the W magnetic moment and $Q_W = -e(1 + \Delta\kappa_\gamma - \lambda_\gamma)/M_W^2$ its quadrupole moment [§]. $(1 + \Delta\mathbf{g}_1^Z)$ can be interpreted as the charge the “Z sees” in the W . Note that the λ terms only involve the field strength, therefore they predominantly affect the production/interaction of transverse W ’s, in other words they do not usefully probe the $S\mathcal{B}$ sector I am keen to talk about here.

Pursuing this observation a little further one can easily describe the distinctive effects the other terms have on different reactions and the reason that some are found to be much better constrained in some reactions than others. First, wherever you look, the λ ’s live in a world on their own, in the “transverse world”. If their effect is found to

[§]The deviations from the minimal gauge value are understood to be evaluated at $k^2 = 0$.

increase dramatically with energy this is due to the fact that these are higher order in the energy expansion (many-derivative operators). The other couplings can also grow with energy if a maximum number of longitudinals are involved, the latter provide an enhanced strength due to the fact that the leading term of the longitudinal polarisation is $\propto \sqrt{s}/M_W$. This enhanced strength does not originate from the field strength!. For instance, in $e^+e^- \rightarrow W^+W^-$, g_1^Z produces one W longitudinal and one transverse: since the produced W come, one from the field strength the other from the “4-potential” (longitudinal) whereas the the κ terms produce two longitudinals and will therefore be better constrained in $e^+e^- \rightarrow W^+W^-$. The situation is reversed in the case of $pp \rightarrow WZ$. This also tells us how one may disentangle between different origins, the reconstruction of the W and Z polarisation is crucial. I have illustrated this in fig. 1, where I have reserved the **thick** arrows for the “important” directions:

Figure 1: *The effect of the phenomenological parameters on the vector boson pair production.*



There are some limits on these couplings from CDF/D0[5] extracted from the study of WZ , WW and $W\gamma$ production: $-2.3 < \Delta\kappa_\gamma < 2.2$; $-0.7 < \lambda_\gamma < 0.7$ while a constrained global fit with $\lambda_\gamma = \lambda_Z$, $\kappa_\gamma = \kappa_Z (g_1^Z = 1)$ gives $-0.9 < \Delta\kappa_V < 1.$; $-0.5 < \lambda_V < 0.5$. I would like to argue that these values are too large to be meaningful. These are too large in the sense that they can hardly be considered as precision measurements, a far cry from the precision that one has obtained on the vector-fermion couplings at LEP1! In the case of the Tevatron and W self-couplings one is talking about deviations of order 100%!

Talking about the LEP1 data, with this year’s statistics one is now sensitive to the genuine non-Abelian radiative corrections and therefore to the presence of the tri-linear (and quadrilinear) couplings [6]. Even so, the data gives no clear information about the presence of the Higgs. In my view this should be taken as very strong evidence for the $SU(2) \times U(1)$ local gauge symmetry or more precisely that the higher order terms that

may correct (2) must *naturally* be small. On the other hand the \mathcal{SB} sector apart from the mass terms still keeps its secret.

It is worth stressing again, contrary to the fierce attack [7] that the above HPZH Lagrangian (eqt. (5)) is not locally gauge invariant and leads to trouble at the quantum level, that as the lengthy introduction has shown all the above operators can be made gauge invariant, by unravelling and making explicit the compensating Goldstone fields and extra vertices that go with the above. Under this light, the HPZH parametrisation should be considered as being written in a specific gauge and that after this gauge (unitary) has been chosen it is non-sensical to speak of gauge invariance[8]. But of course, it is much much better to keep the full symmetry so that one can apply the Lagrangian to any situation and in any frame. There is another benefit in doing so. If the scale of new physics is far enough compared to the typical energy where the experiment is being carried out[¶], then one should only include the first operators in the energy expansion, beyond those of the \mathcal{SM} . Doing so will maintain some constraints on the parameters $\lambda, g_1^Z, \Delta\kappa$. These constraints will of course be lost if you allow higher and higher order operators or allow strong breaking of custodial symmetry, in both cases rendering the situation chaotic while LEP1 shows an incredible regularity. It is highly improbable that the order and symmetry is perturbed so badly.

So what are these operators that describe the self-couplings when one restricts one-self to next-to-leading operators by exploiting the $SU(2) \times U(1)$ and the custodial symmetry? and how are they mapped on the HPZH phenomenological parameters? These are given in Table 1. for the linear [9, 7] as well as the non-linear realisation[1, 10] to bring out some distinctive features about the two approaches:

By going to the physical gauge, one recovers the phenomenological parameters with the *constraints*:

$$\begin{aligned}
\Delta\kappa_\gamma &= \frac{e^2}{s_w^2} \frac{v^2}{4\Lambda^2} (\epsilon_W + \epsilon_B) = \frac{e^2}{s_w^2} \frac{1}{32\pi^2} (L_{9L} + L_{9R}) \\
\Delta\kappa_Z &= \frac{e^2}{s_w^2} \frac{v^2}{4\Lambda^2} (\epsilon_W - \frac{s_w^2}{c_w^2} \epsilon_B) = \frac{e^2}{s_w^2} \frac{1}{32\pi^2} \left(L_{9L} - \frac{s_w^2}{c_w^2} L_{9R} \right) \\
\Delta g_1^Z &= \frac{e^2}{s_w^2} \frac{v^2}{4\Lambda^2} \left(\frac{\epsilon_W}{c_w^2} \right) = \frac{e^2}{s_w^2} \frac{1}{32\pi^2} \left(\frac{L_{9L}}{c_w^2} \right) \\
\lambda_\gamma &= \lambda_Z = \left(\frac{e^2}{s_w^2} \right) L_\lambda \frac{M_W^2}{\Lambda^2}
\end{aligned} \tag{6}$$

Catch 22:

[¶]If this is not the case then we should see new particles or at least detect their tails.

Table 1: *The Next-to-leading Operators describing the W Self-Interactions which do not contribute to the 2-point function.*

Linear Realization , Light Higgs	Non Linear-Realization , No Higgs
$\mathcal{L}_B = ig' \frac{\epsilon_B}{\Lambda^2} (\mathcal{D}_\mu \Phi)^\dagger B^{\mu\nu} \mathcal{D}_\nu \Phi$	$\mathcal{L}_{9R} = -ig' \frac{L_{9R}}{16\pi^2} \text{Tr}(\mathbf{B}^{\mu\nu} \mathcal{D}_\mu \Sigma^\dagger \mathcal{D}_\nu \Sigma)$
$\mathcal{L}_W = ig' \frac{\epsilon_w}{\Lambda^2} (\mathcal{D}_\mu \Phi)^\dagger (2 \times \mathbf{W}^{\mu\nu}) (\mathcal{D}_\nu \Phi)$	$\mathcal{L}_{9L} = -ig' \frac{L_{9L}}{16\pi^2} \text{Tr}(\mathbf{W}^{\mu\nu} \mathcal{D}_\mu \Sigma \mathcal{D}_\nu \Sigma^\dagger)$
$\mathcal{L}_\lambda = \frac{2i}{3} \frac{L_\lambda}{\Lambda^2} g^3 \text{Tr}(\mathbf{W}_{\mu\nu} \mathbf{W}^{\nu\rho} \mathbf{W}^\mu{}_\rho)$	-----
-----	$\mathcal{L}_1 = \frac{L_1}{16\pi^2} \left(\text{Tr}(D^\mu \Sigma^\dagger D_\mu \Sigma) \right)^2 \equiv \frac{L_1}{16\pi^2} \mathcal{O}_1$
-----	$\mathcal{L}_2 = \frac{L_2}{16\pi^2} \left(\text{Tr}(D^\mu \Sigma^\dagger D_\nu \Sigma) \right)^2 \equiv \frac{L_2}{16\pi^2} \mathcal{O}_2$

Aren't there other operators with the same symmetries that appear at the *same level* in the hierarchy and would therefore be as likely?

Answer: YES. And this is an upsetting conceptual problem. On the basis of the above symmetries, one can not help it, but there are other operators which contribute to the tri-linear couplings and have a part which corresponds to bi-linear anomalous W self-couplings. Because of the latter and of the unsurpassed precision of LEP1, these operators are already very much *unambiguously* constrained. Examples of such annoying operators in the two approaches are

$$\begin{aligned} \mathcal{L}_{WB} &= gg' \frac{\epsilon_{WB}}{\Lambda^2} (\Phi^\dagger \times \mathbf{W}^{\mu\nu} \Phi) B_{\mu\nu} \\ \mathcal{L}_{10} &= gg' \frac{L_{10}}{16\pi^2} \text{Tr}(\mathbf{B}^{\mu\nu} \Sigma^\dagger \mathbf{W}^{\mu\nu} \Sigma) \longrightarrow L_{10} = -\pi S \simeq \frac{4\pi s_W}{\alpha} \epsilon_3 \end{aligned} \quad (7)$$

For example, current limits from LEP1 indicate that $-1.4 < L_{10} < 2$. This is really small, so small that if the other L_i 's, say, were of this order it would be extremely difficult to see any effect at the next colliders. So why should the still not-yet-tested operators be much larger? This is the naturalness argument which in my view is *the* essential point of [7]. One can try hard to find models with no contributions to L_{10} . But the solutions are either not very appealing or one has to accept that this quantity can not be calculated reliably in the context of non-perturbative models (in the non-linear approach). Of course, there is also the easy escape that we have not been ingenious enough....

To continue with my talk I will assume that $L_{10} \sim 0$ can be neglected compared to the other operators. This said, I will not completely ignore this limit and the message that LEP1 is giving us, especially that various arguments about the "natural order of

magnitude”^{||} for these operators should force one to consider a limit extracted from future experiment to be meaningful if $|L_i| < \sim 10$ (10 is really generous..). This translates into $\Delta\kappa, \Delta g_Z^1 < \sim 10^{-2}$. Note that the present Tevatron limits if they were to be written in terms of L_9 give $L_9 \sim 10^3$!!!.

With this caveat about L_{10} and the like, let us see how the 2 opposite assumptions about the lightness of the Higgs differ in their most probable effect on the W self-couplings. First, the tri-linear coupling λ is relegated to higher orders in the heavy Higgs limit (less likely). This is as expected: transverse modes are not really an issue here. The main difference is that with a heavy Higgs, genuine quartic couplings contained in $L_{1,2}$ are as likely as the tri-linear and, in fact, when contributing to WW scattering their effect will by far exceed that of the tri-linear. This is because $L_{1,2}$ involve essentially longitudinals. This is another way of arguing that either the Higgs exists or expect to “see something” in WW scattering. Note also that $L_{9L,W,\lambda}$ do give quadri-linear bits but these are imposed by gauge invariance. Note also that $L_{9R,B}$ is not expected to contribute significantly in $pp \rightarrow WZ$ since it has no contribution to Δg_1^Z (see fig. 1). This is confirmed by many analyses.

Going to the physical gauge, the quartic couplings from the chiral approach are

$$\begin{aligned}
\mathcal{L}_{WWV_1V_2}^{SM} = & -e^2 \left\{ \left(A_\mu A^\mu W_\nu^+ W_\nu^- - A^\mu A^\nu W_\mu^+ W_\nu^- \right) \right. \\
& + 2 \frac{c_w}{s_w} \left(1 + \frac{l_{9l}}{c_w^2} \right) \left(A_\mu Z^\mu W_\nu^+ W_\nu^- - \frac{1}{2} A^\mu Z^\nu (W_\mu^+ W_\nu^- + W_\nu^+ W_\mu^-) \right) \\
& + \frac{c_w^2}{s_w^2} \left(1 + \frac{2l_{9l}}{c_w^2} - \frac{l_-}{c_w^4} \right) \left(Z_\mu Z^\mu W_\nu^+ W_\nu^- - Z^\mu Z^\nu W_\mu^+ W_\nu^- \right) \\
& + \frac{1}{2s_w^2} \left(1 + 2l_{9l} - l_- \right) \left(W^{+\mu} W_\mu^- W^{+\nu} W_\nu^- - W^{+\mu} W_\mu^+ W^{-\nu} W_\nu^- \right) \\
& - \frac{l_+}{2s_w^2} \left((3W^{+\mu} W_\mu^- W^{+\nu} W_\nu^- + W^{+\mu} W_\mu^+ W^{-\nu} W_\nu^-) \right. \\
& \left. \left. + \frac{2}{c_w^2} \left(Z_\mu Z^\mu W_\nu^+ W_\nu^- + Z^\mu Z^\nu W_\mu^+ W_\nu^- \right) + \frac{1}{c_w^4} Z_\mu Z^\mu Z_\nu Z^\nu \right) \right\} \\
& \text{with } l_{9l} = \frac{e^2}{32\pi^2 s_w^2} L_{9L} ; \quad l_\pm = \frac{e^2}{32\pi^2 s_w^2} (L_1 \pm L_2)
\end{aligned} \tag{8}$$

Note that the genuine trilinear L_{9L} gives structures analogous to the \mathcal{SM} . The two photon couplings (at this order) are untouched by anomalies.

^{||}Refer to the talk of Wudka.

2 Future Experimental Tests

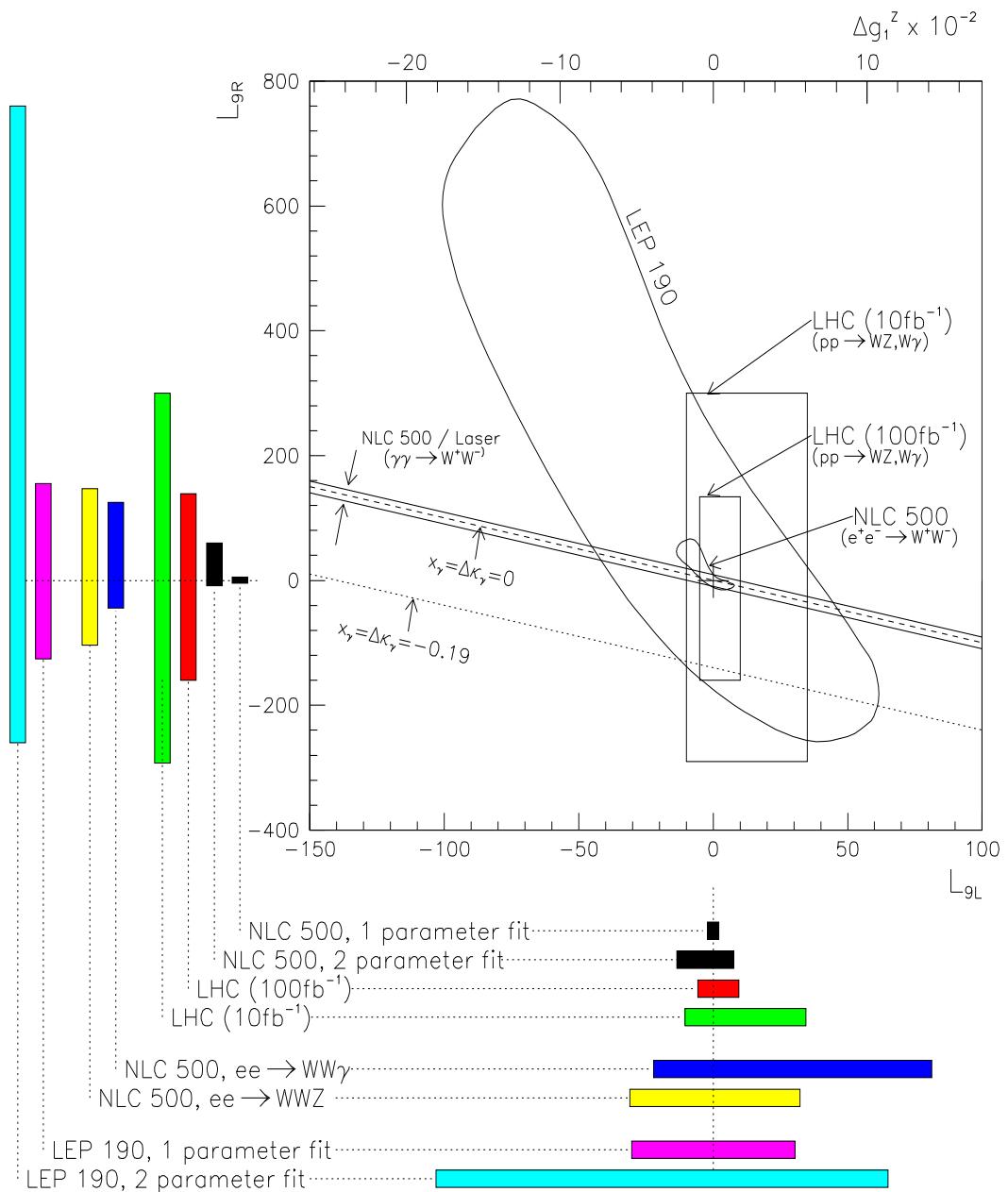
With the order of magnitude on the L_i that I have set as a meaningful benchmark, one should realise that to extract such (likely) small numbers one needs to know the \mathcal{SM} cross sections with a precision of the order of 1% or better. This calls for the need to include the radiative corrections especially the initial state radiation. Moreover one should try to extract as much information from the W and Z samples: reconstruct the helicities, the angular distributions and correlations of the decay products. These criteria mean precision measurements and therefore we expect e^+e^- machines to have a clear advantage assuming that they have enough energy. Nonetheless, it is instructive to refer to fig. 1 to see that pp machines could be complementary.

In the following, one should keep in mind that all the extracted limits fall well within the unitarity limits. I only discuss the description in terms of “anomalous couplings” below an effective cms energy of a VV system $\sim 4\pi v \sim 3 - 4\text{TeV}$, without the inclusions of resonances**. Moreover, I will not discuss the situation when parameters are dressed with energy dependent form factors or any other scheme of unitarisation that introduces more model dependence on the extraction of the limits. For reasons of space I will not go into the details of how the various operators are looked for in various processes and different machines ($pp, e^+e^-, \gamma\gamma$) but refer to a summary I have given elsewhere[1]. I will, however, update some of the results and summarise them in the comparative figure that gives the limits on the genuine tri-linear couplings. These limits are given in terms of the chiral Lagrangian parameters $L_{9L,R}$ or equivalently using (6) in terms of $L_{B,W}$. They can also be re-interpreted in terms of the more usual κ_V, g_1^Z with the constraint given by (6), in which case the L_{9L} axis is directly proportional to Δg_Z^1 .

The limits from pp that I have given in[1] are obsolete (in the present updated version pp means LHC with two settings for the luminosity 10 and 100fb^{-1}). The new limits are based on a very careful study[11] that includes the very important effect of the QCD NLO corrections as well as implementing the full spin correlations for the most interesting channel $pp \rightarrow WZ$. WW production with $W \rightarrow jets$ production is fraught with a huge QCD background, while the leptonic mode is extremely difficult to reconstruct due to the 2 missing neutrinos. The NLO corrections for WZ production are huge, especially in precisely the regions where the anomalous are expected to show up[11]. For instance, high p_T^Z . In the inclusive cross section this is mainly due to, first, the importance of the subprocess $q_1g \rightarrow Zq_1$ (large gluon density at the LHC) followed by the “splitting” of the

**At this conference, this has been discussed by Roberto Casalbuoni and Kingman Cheung.

Figure 2: Comparison between the expected bounds on the two-parameter space $(L_{9L}, L_{9R}) \equiv (L_W, L_B) \equiv (\Delta g_1^Z, \Delta \kappa_\gamma)$ (see text for the conversions) at the NLC500 (with no initial polarisation), LHC and LEP2. The NLC bounds are from $e^+e^- \rightarrow W^+W^-$, $W^+W^-\gamma$, W^+W^-Z (for the latter these are one-parameter fits) and $\gamma\gamma \rightarrow W^+W^-$. The LHC bounds are from $pp \rightarrow WZ$. Limits from a single parameter fit are also shown (“bars”).

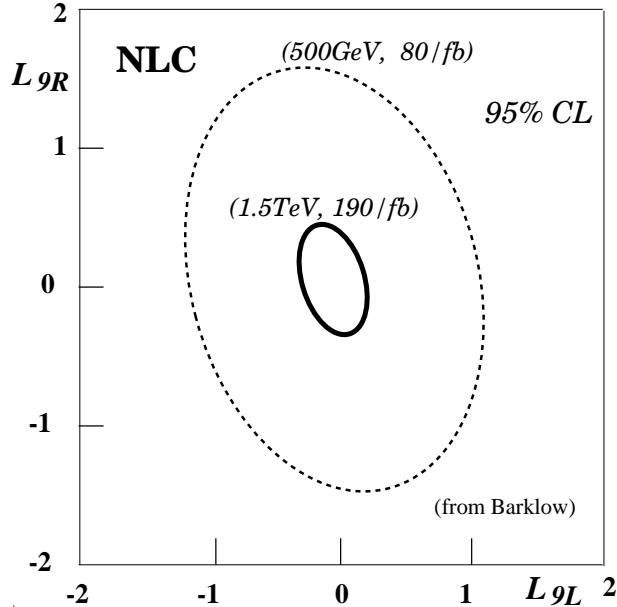


quark q_1 into W . The probability for this splitting increases with the p_T of the quark (or Z): $\text{Prob}(q_1 \rightarrow q_2 W) \sim \alpha_w / 4\pi \ln^2(p_T^2/M_w^2)$. To reduce this effect one [11] has to define an exclusive cross section that should be as close to the LO WZ cross section as possible by cutting on the extra high p_T quark (dismiss any jet with $p_T^{\text{jet}} > 50\text{GeV}$, $|\eta_{\text{jet}}| < 3$). This defines a NLO $WZ + \text{"0jet"}$ cross section which is stable against variations in the choice of the Q^2 but which nonetheless can be off by as much as 20% from the prediction of Born \mathcal{SM} result. The anomalous parameters are included one by one in the form of the HPZH parameterisation. It is indeed found, as expected from the general arguments that I exposed above, that Δg_Z^1 is much better constrained than $\Delta \kappa_Z$. I have thus reinterpreted the results in the chiral Lagrangian approach approximating the effect of L_{9L} as being dominantly due to Δg_Z^1 while I blamed the bad limit on $\Delta \kappa_Z$ on L_{9R} .

For the case of e^+e^- at high energies, the comparative figure shows the adaptation of the *BM2*[12] results. These are based on a very powerful fitting procedure that aims at reconstructing 8 observables which are combinations of density matrices. Simulations performed for LEP2 energies by experimentalists[13] have shown that we can somehow improve on these limits. The main missing ingredient that may change these results is, once again, the effect of radiative corrections. Notably, bremsstrahlung and beamstrahlung were not taken into account. It is now mandatory to include these corrections, for a review see[14]. As in the case of pp , initial state radiation drastically affects some of the distributions that, at tree-level, seem to be good New Physics discriminators. For instance, initial state radiation is responsible for the boost effect that redistributes phase space: this leads to the migration of the forward W into the backward region and results in a large correction in the backward region. Precisely the region where one would have hoped to see any s-channel effect more clearly. Second, if one reconstructs the polarisation of the W without taking into account the energy loss, one may “mistag” a transverse W for a longitudinal, thereby introducing a huge correction in the small tree-level longitudinal cross sections, which again is particularly sensitive to New Physics. Cuts must be included. With the near advent of LEP2, there is now the discussion[15] whether an analysis based on the resonant diagrams is enough. It is found that non resonating (non genuinely WW) 4-fermion states are not negligible. Probably, it is best to cut on the non-resonant diagrams by double mass constraints (etc..) at the expense of reducing the event sample, rather than working with a “mixed” final state.

Very recently Barklow[16] has reanalysed the operators $L_{9L,9R}$ by considering the correlated 4-fermion- WW five-fold angular distributions and including NLC luminosity spectra as well as considering the effect of initial polarisation. The latter, as is known, can easily

Figure 3: *Limits on $(L_{9L} - L_{9R})$ in e^+e^- including ISR and beam polarisation.*



isolate the the s-channel WWV . His analysis at 500GeV assumes a luminosity of 80fb^{-1} , which is much larger than what has been assumed in the similar study of *BM2* (10fb^{-1}). However, since the sensitivity to the anomalous goes like $\sim \sqrt{\mathcal{L}}$ this confirms the *BM2* results and hints that although the inclusion of the luminosity spectra makes the analysis more complicated it, fortunately, does not critically degrade the sensitivity to the anomalous. Anyway, with this luminosity the results are fascinating, one can be sensitive to values as low as 1-2 for the parameters L_9 . This is really precision measurement. Moreover, in future e^+e^- linacs one also hopes to have a $\gamma\gamma$ version. A new analysis[17] shows that in combination with the e^+e^- mode, the $\gamma\gamma$ mode can help put much stronger limits on the parameter space of the anomalous (see fig. 2).

In conclusion, it is clear that already with a 500GeV e^+e^- collider combined with a good integrated luminosity of about 80fb^{-1} one can reach a precision, on the parameters that probe $\mathcal{S}\mathcal{B}$ in the genuine tri-linear WWV couplings, of the same order as what we can be achieved with LEP1 on the two-point vertices. To reach higher precision and critically probe $\mathcal{S}\mathcal{B}$ one needs to go to TeV machines, as fig. 3 shows for the tri-linear $L_{9L} - L_{9R}$. In fact, at an effective WW invariant masses of order the TeV, $\mathcal{S}\mathcal{B}$ (especially in scalar-dominated models) is best probed through the genuine quartic couplings in WW scattering or even perhaps in WWZ, ZZZ production (that are poorly constrained at 500GeV). LHC could also address this particular issue but one needs dedicated careful simulations to see whether any signal could be extracted in the pp environment. In this

regime there is also the fascinating aspect of W interaction that I have not discussed and which is the appearance of strong resonances. This would reveal another alternative to the \mathcal{SM} description of the scalar sector.

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